A Top Quark Soliton and its Anomalous Chromomagnetic Moment

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We show that under the assumption of dynamical symmetry breaking of electro weak interactions by a top quark condensate, motivated by the Top Mode Standard Model, the top quark in this effective theory can be considered then as chiral color soliton (qualiton). This is realized in an effective four-fermion interaction with chiral $SU(3)_c$ as well as $SU(2)_L \otimes U_Y(1)$ symmetry. In the pure top sector the qualiton consists of a top valence quark and a Dirac sea of top and anti-top quark coupled to a color octet of Goldstone pions. The mass spectra, isoscalar quadratic radii and the anomalous chromomagnetic moment due to a non-trivial color form factor are calculated with zero and finite current top masses and effects at the Hadron Colliders are discussed. The anomalous chromomagnetic moment turns out to have a value consistent with the top production rates of the D0- and CDF-measurements.

I. INTRODUCTION

That the mechanism for symmetry breakdown in the electro weak theory of the standard model (SM) is maybe of dynamical origin was first noted by Nambu [1] and Miransky et al. [2] and is presently under intensive discussion [3–7]. The clue is that the recently measured top quark mass of CDF Collaboration [8] $m_t = 176 \pm 8 \pm 10\text{GeV}$ ($\sigma_p = 6.8^{+3.6}_{-2.4} pb$) and D0 Collaboration [9] $m_t = 199^{+19}_{-21} \pm 22\text{GeV}$ ($\sigma_p = 6.4 \pm 2.2 pb$) turn out to be correct and the top mass is of the order of the Fermi scale $v = \left(\sqrt{2}G_F\right)^{-1/2} = 246\text{GeV}$, the scale of electro weak symmetry breaking. Therefore if there is any new physics in the electro weak theory beyond this mass scale, the top quark should be an ideal probe for theses effects which are almost invisible for the lighter quarks.

Also, up to now there is no evidence for any process with Higgs boson contribution at Born level [10]. Furthermore if the D0 Measurement turns out to be correct, the bounds on the Higgs will be in the TeV regime [10], whereas baryon number violating processes in the minimal Standard model (cf. Diakonov et al. [11]) presumably restrict m_H to be smaller than 80GeV. Having this in mind it is not too speculative to think of new physics beyond the electroweak theory.

For this purpose we want to consider a model Lagrangian on the electroweak scale hoping that it contains the relevant degrees of freedom of the real unified theory at some higher energy scale. The low energy behaviour should of course coincide with the standard electro weak theory [12]. One of the candidates for such an effective theory is the so called *Top Mode Standard Model* [1,2] where the dynamical symmetry breaking is performed via a top quark condensate in a BCS oder Nambu–Jona-Lasinio like theory [13]. Because the top quark Yukawa coupling is expected to be of the order of one, the situation is similar to the nucleon constituent quarks in QCD. This makes the assumption of a dynamical symmetry breaking reasonable, which gives a mass to the top quark, comparable to the up- and down-constituent quarks in QCD, which are known to get their masses dynamically.

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Recently bounds on the radii for the up, down, strange, charm and bottom quarks have been reported [14] and they are actually very small, i.e. $r < 10^{-5} fm$. So they can be considered as point like with no non-trivial formfactor and therefore vanishing anomalous magnetic moment. For the top quark the situation is probably different due to its large mass which is very much larger than the masses of the other fermions.

It was noted by Brodsky and Drell [15] and recently by Brodsky and Schlumpf [16] that the anomalous magnetic moment is presumably linear in the size and mass of the composite particle. If such an picture also applies to the top quark, the large mass in conjunction with a finite radius ($\simeq 10^{-3}$fm) should give a measurable effect.

But actually it was proposed already that a discrepancy between SM cross section for top production [17] and first CDF measurements [18] can be explained by a non-trivial formfactor for the top (see Ref. [19] for a short review). Therefore we will investigate here the idea proposed by Atwood, Kagan and Rizzo [20] that already a small anomalous chromomagnetic moment for the top quark can enhance the $\bar{t}t$ production cross section [1]. We will find, that our model gives an anomalous magnetic moment, which has the correct magnitude to explain the production rate at the Tevatron.

The formalism that we will use to calculate the anomalous chromomagnetic moment for the top quark is an effective Lagrangian with BCS or Nambu–Jona-Lasinio (NJL) type mechanism for dynamical symmetry breaking. The idea is that at some high energy scale Λ the SM contains only the usual quark, lepton and gauge boson degrees of freedom, but no fundamental Higgs scalars [4]. The dynamical symmetry breaking in the electroweak sector is then triggered by the condensate of the top.

The formalism of solitons having non-integer baryon number was suggested by Kaplan [23] and later investigated by Gomelski and Karliner [24] and Keaton [25] in the framework of the *Skyrme model*, where the baryon number is induced topologically via the winding number of the chiral profile. The constituent quarks obtained their masses via spontaneously broken chiral SU(3) color symmetry. In the early work [23] Kaplan wants to calculate the properties of the three lightest constituent quarks in order to deduce finally the properties of the hyperons which should consist of these constituent quarks. However, apart from the problem of having too large constituent quark masses after the quantization, the problem of treating the baryons as bound states of these constituent quarks was never solved.

Recently the formalism of fractional baryon number was refreshed by Zhang [3] for the top quark in the electro weak theory.

In the Top Mode Standard Model (TMSM) (see e.g. [4–6,26,27]) the chiral color symmetry is used to trigger the formation of a top quark condensate via a local *four point fermion* interaction Lagrangian. However the original formulation suffered from some fine-tuning problems and an enormous ultraviolet cutoff.

We want to make these ideas more explicit by considering the fermion determinant of an effective top quark interaction Lagrangian, suggested by Lindner and Ross [6], and look for selfconsistent solitonic solutions of the classical Euler Lagrange equations. In contrast to the treatment of this Lagrangian in the Top Mode Standard Model, where a top quark condensate leads to symmetry breaking and the top mass is given by a effective potential or the so called gap equation, we obtain in our approach a localized solution of the corresponding solitonic equations of motion, the top quark soliton (qualiton). The top quark constituent mass will be – as we shall see later – the free parameter of our model. The picture that emerge is that of a top valence quark carrying around with it a cloud of Goldstone bosons interacting with the polarized Dirac sea of top and antitop quarks. The fractional baryon number 1/3 of this soliton is then non-topological and induced by the explicit occupation of the discrete valence level.

Adopting the semiclassical quantization scheme [28] in order to project out of the soliton a state with definite color and spin, we calculate the mass of the top quark in the fundamental representation and the low lying resonances. Furthermore we look at observables like the isoscalar quadratic radius and the chromomagnetic moment. The latter quantity is of special interest since it should affect top production rates and distributions at Hadron Colliders [29,30].

The organization of the paper is as follows. In sect. II we determine the parameters of our Lagrangian by fixing decay constant of the color octet of Goldstone bosons. In sect. III we derive the classical and quantum mass of the top states. In sect IV we calculate the anomalous magnetic moment of the top ground state and

[1]See also e.g. [21,22] for a discussion of non-standard top-quark couplings.

in sect. V we review the numerical results. In sec. VI we finally summarize our results. The appendix A is devoted to a careful definition of the quantization of fractional baryon numbers.

II. THE LAGRANGIAN

In order to account for the physics beyond the Standard Model several effective Lagrangians [1,2,4] have been proposed all having in common the spontaneous breaking of electro weak symmetries by a top quark condensate. Since we are in this first step neither interested in the various gauge fields itselves nor in the detailed mechanism of spontaneous breaking of $SU(2)_L \otimes U_Y(1)$ our starting point is, according to Lindner and Ross [6], the effective interaction (at momenta below a few TeV)

$$\mathcal{L}_{int} = -G \quad \bar{L}_{\alpha,a} \delta^c_{\alpha\beta} \gamma^\mu_{ab} L_{\beta,b} \quad \bar{t}_{R,\gamma,c} \delta^c_{\gamma\delta} \gamma^\mu_{cd} t_{R,\delta,d} \tag{1}$$

where $L = (t_L, b_L)$ is the lefthanded fermion doublet of the electroweak theory and greek letters imply summation over color indices and roman letters over spin indices. As a trivial check, this model has color chiral $SU(3)_L \otimes SU(3)_R$ symmetry as well as $SU(2)_L \otimes U_Y(1)$. Using Fierz identities such as

$$(\gamma^\mu)_{ab}(\gamma^\mu)_{cd} = (K^\alpha)_{ad}(K^\alpha)_{cb},\ K^\alpha = (\mathbf{1}, i\gamma_5, \frac{i}{\sqrt{2}}\gamma_\mu, \frac{i}{\sqrt{2}}\gamma_\mu\gamma_5) \tag{2}$$

and

$$\delta^c_{\alpha\beta}\delta^c_{\gamma\delta} = \frac{1}{2}\sum_{a=0}^{8}(\lambda^a)_{\alpha\delta}(\lambda^a)_{\gamma\beta} \tag{3}$$

the interaction eq. (1) can be rewritten as

$$\mathcal{L}_{int} = G\frac{1}{2}\left[\bar{L}\mathbf{1}\lambda^a t\bar{t}\mathbf{1}\lambda^a L + \bar{L}i\gamma_5\lambda^a t\bar{t}i\gamma_5\lambda^a L\right]. \tag{4}$$

Using $\gamma_5 t = t$ and $\gamma_5 L = -L$ this reduces to

$$\mathcal{L}_{int} = G\left[\bar{L}\mathbf{1}\lambda^a t\bar{t}\mathbf{1}\lambda^a L\right] \tag{5}$$

and with the non-chiral decompositions $L = \frac{1}{2}(1-\gamma_5)Q$ and $t = \frac{1}{2}(1+\gamma_5)q$ we obtain

$$\begin{aligned}\mathcal{L}_{int} &= \frac{G}{4}\left[\bar{Q}\lambda^a q + \bar{Q}\gamma_5\lambda^a q\right]\left[\bar{q}\lambda^a Q - \bar{q}\gamma_5\lambda^a Q\right] \\ &= G\left[\bar{Q}\lambda^a q\bar{q}\lambda^a Q - \bar{Q}\gamma_5\lambda^a q\bar{q}\gamma_5\lambda^a Q + \bar{Q}\gamma_5\lambda^a q\bar{q}\lambda^a Q - \bar{Q}\lambda^a q\bar{q}\gamma_5\lambda^a Q\right].\end{aligned} \tag{6}$$

Written in terms of bottom and top fields it is given by

$$\mathcal{L}^{tt}_{int} = \frac{G}{4}\left[\bar{t}\lambda^a t\bar{t}\lambda^a t - \bar{t}\gamma_5\lambda^a t\bar{t}\gamma_5\lambda^a t\right] \qquad \text{and} \tag{7}$$

$$\mathcal{L}^{tb}_{int} = \frac{G}{4}\left[\bar{b}\lambda^a t\bar{t}\lambda^a b - \bar{b}\gamma_5\lambda^a t\bar{t}\gamma_5\lambda^a b - \bar{b}\lambda^a t\bar{t}\gamma_5\lambda^a b + \bar{b}\gamma_5\lambda^a t\bar{t}\lambda^a b\right]. \tag{8}$$

Ignoring the bottom quark degree of freedom at this point, because their Yukawa coupling is much lower than that of the top mass, the usual bosonization procedure [31] leads to

$$\mathcal{L}^{tt}_{int} = f_t\ \bar{t}\left(\bar{\sigma} + i\bar{\pi}\gamma_5\right)t + \frac{\chi^2}{4}\mathrm{tr}\left(\bar{\sigma}^2 + \bar{\pi}^2\right) \tag{9}$$

where $G/2 = f_t^2/\chi^2$ and $\bar{\sigma}$ and $\bar{\pi}$ are real SU(3) matrix fields. These are in the adjoint representation of $SU(3)_c$. Integrating out the quarks gives the one loop effective action

$$S^{tt}_{eff} = -\mathrm{Sp}\ \log\left(-i\partial\!\!\!/ + f_t(\bar\sigma + i\bar\pi\gamma_5)\right) + \frac{\chi^2}{4}\int d^4x \mathrm{tr}\left(\bar\sigma^2 + \bar\pi^2\right) \tag{10}$$

A stationary phase condition for the vacuum leads to a gap equation of the spontaneously broken phase with $\bar\sigma_v \neq 0$ and

$$\chi^2 = 8 f_t^2 I_1(M_{con}), \tag{11}$$

where $M_{con} = f_t \bar\sigma_v$ is a diagonal constituent quark mass matrix for the top. As we will see later, this is not the physical mass of the top but rather the classical one which will be modified in the presence of the non-homogeneous Dirac sea. In eq. (11)

$$I_n(M) = \frac{1}{16\pi^2}\int \frac{du}{u^{3-n}} e^{-uM^2}\phi(u), \qquad \phi(u) = \sum_{i=1}^{n} \theta(1 - \frac{1}{\Lambda_i^2}) \tag{12}$$

are proper-time regularized [32] Feynman integrals. We gave in eq. (12) a rather general form of the cutoff function as a sum over n step like cutoff functions [33]. However for our present purpose we will take $n = 1$, which is sufficient to render the model finite. More cutoffs could be used to fix condensates or current masses [33]. In choosing the vacuum state $\bar\sigma_v \neq 0$ in eq. (11) the original $SU(3)_R \otimes SU(3)_L$ symmetry is spontaneously broken down to $SU(3)_V$. As a result the pion-like fields $\bar\pi$ are the Goldstone bosons of the spontaneously broken symmetry. The usual wave-function renormalization condition $Z_\pi = 1$ [31,34,35] gives

$$4 M_{con}^2 I_2(M_{con}) = f_\pi^2 \tag{13}$$

and determines the value of the cutoff Λ (12) in terms of the constituent mass M_{con}. Thus for a given top-pion decay constant f_π, the top-constituent mass M_{con} is the only free parameter of the model. The top-pion decay constant is of course not yet known but it is expected to be of the order of the Fermi scale. We choose a central value of

$$f_\pi = f_\pi^{tt} \simeq 30\mathrm{GeV} \tag{14}$$

which is of the order of $f_\pi^{tb} \simeq 50$GeV proposed by Hill [36] in some related technicolor model. However this fixing is by no means unique and we also studied variations of this central value in Tab. (II) and Tab. (III). However – as we shall see later – the anomalous magnetic moment will turn out to be independent of f_π.

III. SOLITONS

In order to describe a baryon number $\frac{1}{3}$ system from the effective action eq. (10), we have to consider meson fields $\bar\sigma$ and $\bar\pi$ different from the vacuum configuration. Because we are interested in static properties of the top, we restrict ourselves to time-independent meson fields.

Considering first the SU(2) version of the model, one has to make the replacement

$$\bar\sigma + i\bar\pi\gamma_5 \longrightarrow f_\pi \xi = f_\pi e^{i\theta\hat{x}\vec\tau\gamma_5}. \tag{15}$$

This combines two things. First the usual hedgehog ansatz for the quark fields allows only the isoscalar scalar and isovector pseudoscalar fields to have non-trivial equation of motion. Second a non-linear constraint for $\bar\sigma_0$ and $\bar\pi$ has to be used in order to get stable solitonic solutions [37]. So the meson field configuration is given by the chiral angle θ, which can be obtained from the effective action eq. (10) by solving the corresponding equations of motion [38,39,54].

By defining a **Giantspin** $\vec P = \vec j + \vec\lambda$ analogous to the usual Grandspin $\vec G = \vec j + \vec\tau$, the hamiltonian

$$H = \frac{\vec\alpha\cdot\vec\nabla}{i} + \beta M_{con}\left(\cos\theta(r) + i\gamma_5 \sin\theta(r)\frac{\vec x}{r}\cdot\vec\lambda\right) \tag{16}$$

commutes with Giantspin P^2, its z-component P_z as well as the parity Π and these four quantities form a complete set of commuting operators. The eigenstates $|n\rangle$ of H:

$$H|n\rangle = E_n|n\rangle \qquad \langle \vec{r} \mid n\rangle = \phi_n(\vec{r}) \tag{17}$$

can be characterized by four quantum numbers E_n, P, P_z and $(-)^l$, where P_z is degenerated. The eigenvalue problem eq. (17) will now be solved numerically following [40] by putting the system in a large, finite box and demanding appropriate boundary conditions for the radial part of the eigenfunctions at the end of the box. So we obtain a discrete basis, which belongs to a given giant spin P and parity Π. With a numerical cutoff the basis became finite, so that a numerical diagonalization is possible. Note that this numerical cutoff has nothing to do with the physical UV cutoff Λ mentioned above.

The equations of motion for our system are given by stationary points of the effective chiral action. These are solvable via a standard selfconsistent procedure as it is known from Hartree-Fock calculations in nuclear physics [41,39]. We start with a reasonably chosen profile $\theta(r)$, diagonalize the H of eg. (16) as described above, and obtain eigenfunctions $\phi_n(\vec{r})$ and eigenvalues E_n, which lead to a new profile through the equations of motion. This selfconsistent procedure is iterated until a reasonable small difference in the profiles is reached.

In general the selfconsistent procedure is performed for a given value of the constituent top mass M_{con} (or equivalently the scaled cutoff Λ/M_{con}, see Tab. (II)). Furthermore explicit symmetry breaking current masses as well as no symmetry breaking current mass have been considered (compare Tab. (III) and Tab. (A)). It turns out,that solitonic solutions exists only if the scaled cutoff Λ lies under a critical value $\Lambda < \Lambda_{\mathrm{cr}} \approx 2$, which corresponds to lower bound of the constituent top mass of $M_{\mathrm{cr}} \approx 180$ GeV for $f_\pi = 30$ GeV. This critical bound is typical for localized solutions of a system of coupled non linear equations [42] (solitonic solutions).

In dependence of the radial size of the profile θ, one finds bound orbitals both in the positive and negative spectrum of eigenstates. Due to the spontaneous symmetry breaking there is a gap of size $2 \cdot M_{\mathrm{con}}$ between these parts of the spectrum and the positive bound state with lowest single particle energy E_n and quantum numbers $P^+ = 0^+$ will be called valence orbit E_{val}. Its energy decreases with increasing profile size, switches sign and gets finally part of the negative spectrum [41]. This valence level, coming from the positive continuum, gets bound and localized by interacting with the negative continuum, which gives us the solitonic solution of the system.

An ansatz for the SU(3) fields is given by using the trivial embedding [43,23] of the subgroup of SU(2) into SU(3). Then the unitary field ξ performs to

$$\xi = \begin{pmatrix} e^{i\theta(r)\hat{x}\vec{\lambda}\gamma_5} & 0 \\ 0 & 1 \end{pmatrix}, \tag{18}$$

This embedding made by Witten for the flavor SU(3) group is distinguished by giving the correct selection rules for the low lying multiplets with definite color and spin.

The hedgehog mean field solution of the classical equations of motion breaks the rotational symmetry of the full theory, i.e. the eigenstates $|n\rangle$ of H do not carry good spin quantum number. To restore this symmetry, we couple the corresponding spin expectation value to the effective chiral action, which turns out to be equivalent to consider the soliton in a rotating system, called cranking method [44,45]. The main idea is to perform an adiabatic rotation of the hedgehog fields with angular velocity $\vec{\Omega}$ and treat the problem perturbatively in Ω. In this expansion there occur rotational terms subleading in Ω. They are not necessarily small due to their – long time overlooked – origin [48–50]: The model is formulated in terms of a fermion determinant (10), so that the operators, which occur in the path integral formulation for matrix elements have to be explicit time-ordered. Supplementary the cranking velocities of the semiclassical quantization turn out to be collective operators, which do not commute with the rotation matrix itself.

The cranking procedure gives us the moment of inertia tensor

$$I_{AB} = \frac{1}{4}\int \frac{d\omega}{2\pi} \mathrm{Tr}_{\gamma\lambda_c} \frac{1}{i\omega + H}\lambda_A \frac{1}{i\omega + H}\lambda_B. \tag{19}$$

Because of the embedding we have

$$I_{AB} = \begin{cases} I_1\delta_{AB} & \mathrm{for} A, B = 1,2,3 \\ I_2\delta_{AB} & \mathrm{for} A, B = 4,5,6,7 \\ \quad 0 & \mathrm{for} A, b = 8 \end{cases} \tag{20}$$

and by defining right generators R_a [46,47]

$$R_a = \begin{cases} I_1\Omega_a & a=1,2,3 \\ I_2\Omega_a & a=4,5,6,7 \\ \frac{1}{2\sqrt{3}} := \frac{\sqrt{3}}{2}Y_R & a=8 \end{cases} \tag{21}$$

we can see that the right hypercharge is restricted to 1/3. This corresponds to multiplet representations with triality 1. The lowest SU(3) representation with unit triality is therefore the triplet **[3]** with spin 1/2. The next one is the color antisextett with spin 1/2, which can be considered as some $t(\bar{t}t)$ excitation. Because we have $N_c = 1$, the baryon number and therefore the right hypercharge Y_R is now 1/3 and the lowest possible representation therefore is the triplet **[3]**, the fundamental representation of SU(3). By expressing the spin expectation value and the rotational contribution of $S_{\text{eff}}[\Omega]$ in terms of the SU(3) rotation matrix we can perform the quantization in the standard way by substituting the coordinates and conjugate momenta through the operators which should fulfill the usual commutator rules. After subtracting the spurious zero mode contributions this leads to

$$M = M_{cl} + \frac{1}{2I_2}C_2(SU(3)) + \left(\frac{1}{2I_1} - \frac{1}{2I_2}\right)C_2(SU(2)) - \frac{3}{8I_1} - \frac{7}{24I_2}, \tag{22}$$

where $C_2(SU(2)_R) = j(j+1) = 3/4$ and $C_2(SU(3)_{R/L}) = 1/3(p^2+q^2+3(p+q)+pq) = 4/3$ for the color triplet **[3]** and spin 1/2 representation (see e.g. [23,3]). So in the lowest representation $(3, \frac{1}{2})$ with $(p=1, q=0)$ the physical top mass M_{top} is simply given by

$$M_{top} = M^{(3,\frac{1}{2})} = M_{cl}. \tag{23}$$

According to our numerical results, to be presented in sect. VI, the next higher representation $(\bar{6}, \frac{1}{2})$ with $(p=0, q=2)$ is shifted by[2]

$$M^{(\bar{6},\frac{1}{2})} - M^{(3,\frac{1}{2})} = \frac{1}{I_2} \simeq 380\text{GeV} \tag{24}$$

for $M_{con} = 190$GeV, $f_\pi = 30$GeV, whereas $(15, \frac{1}{2})$ in the $(p=2, q=1)$ multiplet gives

$$M^{(15,\frac{1}{2})} - M^{(3,\frac{1}{2})} = \frac{2}{I_2} \simeq 760\text{GeV}. \tag{25}$$

The next higher spin multiplet $(15, \frac{3}{2})$ gives

$$M^{(15,\frac{3}{2})} - M^{(3,\frac{1}{2})} = \frac{3}{2I_1} + \frac{1}{2I_2} \simeq 460\text{GeV} \tag{26}$$

which leads to the remarkable fact, that the spin 3/2 multiplet is below a spin 1/2 multiplet.

The moments of inertia in (22) are decomposed into *valence* and *sea* parts $I_k = I_k^{val} + I_k^{sea}$, $k = 1, 2$, according to

$$\begin{aligned} I_1^{val} &= \frac{1}{6}\sum_{i=1}^{3}\sum_{n\neq val}\frac{\langle val \mid \lambda_i \mid n\rangle\langle n \mid \lambda_i \mid val\rangle}{E_{val}-E_n}, \quad i=1,2,3 \\ I_2^{val} &= \frac{1}{8}\sum_{a=4}^{7}\sum_{n\neq val}\frac{\langle val \mid \lambda_a \mid n\rangle\langle n \mid \lambda_a \mid val\rangle}{E_{val}-E_n}, \quad a=4,5,6,7 \end{aligned} \tag{27}$$

And

[2]Note that this shift differs from eq. (4.29) in [23].

$$
\begin{aligned}
I_1^{sea} &= \frac{1}{12}\sum_{i=1}^{3}\sum_{m,n}\langle m \mid \lambda_i \mid n\rangle\langle n \mid \lambda_i \mid m\rangle \mathcal{R}_I(E_n, E_m), \quad i = 1,2,3 \\
I_2^{sea} &= \frac{1}{16}\sum_{a=4}^{7}\sum_{m,n}\langle m \mid \lambda_a \mid n\rangle\langle n \mid \lambda_a \mid m\rangle \mathcal{R}_I(E_n, E_m), \quad a = 4,5,6,7
\end{aligned}
\tag{28}
$$

where $\mathcal{R}_I(E_n, E_m)$ is the usual regularization function for the sea part of moment of inertia [33]:

$$
\begin{aligned}
\mathcal{R}_I(E_n, E_m) = -\frac{1}{2\pi}\int_0^\infty \frac{du}{\sqrt{u}}\sum_{i=1}^{n}\theta\left(u - \frac{1}{\Lambda_i^2}\right)\frac{1}{E_m + E_n} \times \\
\left[E_n \exp(-uE_n^2) + E_m \exp(-uE_m^2) + \frac{\exp(-uE_n^2) - \exp(-uE_m^2)}{u(E_n - E_m)}\right].
\end{aligned}
\tag{29}
$$

IV. CHROMOMAGNETIC MOMENTS

The color octet of gauge fields $G_\mu^a(x)$ couple in the QCD Lagrangian to the top quarks $\mathcal{L}_{ttG} = g\ j_\mu^a(x)\ G_\mu^a(x)$ where $j_\mu^a(x) = \bar{t}(x)\gamma_\mu\lambda^a\frac{1}{2}t(x)$ is the conserved color octet vector current. From this we can define the **chromomagnetic moment** of the top as

$$
\mu_z^c = \frac{g}{2M_{top}}\int d^3x\left(\vec{x} \times \vec{j}^c\right)_z =: \frac{g}{2M_{top}}G_M^c(q^2 = 0)
\tag{30}
$$

where g is the gluon coupling constant, related to α_s in QCD. The general form of the matrix element of the vector current for a localized top quark state can be described by **color formfactors** F_1 and F_2 as:

$$
\sqrt{\frac{p_0 p_0'}{M_{top}^2}}\langle t(p,c,s) \mid j_\mu^a(0) \mid t(p',c,s)\rangle = \bar{t}_{c,s}(p)\left[F_1(q^2)\gamma_\mu + i\frac{\sigma_{\mu\nu}q_\nu}{2M_{top}}F_2(q^2)\right]\lambda^c\frac{1}{2}\ t_{c,s}(p').
\tag{31}
$$

From the color formfactors one can calculate the chromoelectric $G_E^c(q^2)$ and chromomagnetic $G_M^c(q^2)$ formfactors according to

$$
\begin{aligned}
G_E^c(q^2) &= G_E(q^2)L_c = \left[F_1(q^2) + \frac{q^2}{4M_{top}^2}F_2(q^2)\right]L_c \\
G_M^c(q^2) &= G_M(q^2)L_c = \left[F_1(q^2) + F_2(q^2)\right]L_c
\end{aligned}
\tag{32}
$$

where L_c is eigenvalue of the left color SU(3) generator. Now it can be shown that $F_1(0)$ is related to the normalized color charge and therefore $F_1(0) = 1$. Then the *anomalous* chromomagnetic moment $\kappa = F_2(0)$ follows as

$$
\kappa = F_2(0) = G_M(0) - G_E(0) = G_M(0) - 1.
\tag{33}
$$

Remember that for a point like particle $F_1 = 1$ and $F_2 = 0$. Similarly in low energy QCD regime we have for the proton $F_1 = 1$ and $F_2 = 1.79$. Therefore the prediction of a non-zero $\kappa = F_2(0)$ is a result of the color form factor of the top and is absent for point-like quarks[3]. However the reverse is not true. We will show that $\kappa = 0$ and $R \neq 0$ is consistent. Actually we show that $\kappa + 1$ is proportional to $M_{top} \cdot R$ in contrast to κ itself in Ref. [16].

Similar to the axial current formfactor in Ref. [48] within a flavor SU(3) model the chromomagnetic formfactor at $q^2 = 0$ can be written as

[3]Provided that non-universal interactions are neglected [21,22].

$$G_M^c(q^2=0) = \left(A + \frac{B}{I_1} + \frac{C}{I_2}\right) D_{ci} - \frac{D}{I_2} d_{ibb} D_{cb} R_b - \frac{E}{I_1} D_{c8} R_i, \quad i=3 \tag{34}$$

where the coefficients A, B, C, D, E are given by

$$\begin{aligned}
A &= \frac{M_{top}}{6} \sum_n \langle n \mid \gamma_0 \gamma_i \lambda_3 \hat{x}_j \mid n \rangle \epsilon_{i3j} \mathcal{R}_\Sigma(E_n) \\
B &= \frac{M_{top}}{12} \sum_{m,n} \langle n \mid \gamma_0 \gamma_i \lambda_k \hat{x}_j \mid m \rangle \langle m \mid \lambda_l \mid n \rangle \epsilon_{ij3} \epsilon_{kl3} \mathcal{R}_\mathcal{Q}(E_n, E_m) \\
C &= \frac{M_{top}}{12} \sum_{m,n} \langle n \mid \gamma_0 \gamma_i \lambda_a \hat{x}_j \mid m \rangle \langle m \mid \lambda_b \mid n \rangle \epsilon_{ij3} \left(f_{3ab} - \epsilon_{3ab}\right) \mathcal{R}_\mathcal{Q}(E_n, E_m) \\
D &= \frac{M_{top}}{12} \sum_{m,n} \langle n \mid \gamma_0 \gamma_i \lambda_a \hat{x}_j \mid m \rangle \langle m \mid \lambda_a \mid n \rangle \; \epsilon_{ij3} \; d_{3aa} \mathcal{R}_\mathcal{M}(E_n, E_m) \\
E &= \frac{M_{top}}{12} \sum_{m,n} \langle n \mid \gamma_0 \gamma_i \lambda_a \hat{x}_j \mid m \rangle \langle m \mid \lambda_b \mid n \rangle \; \epsilon_{ij3} \; \delta_{a8} \; \delta_{b3} \mathcal{R}_\mathcal{M}(E_n, E_m)
\end{aligned} \tag{35}$$

The regularization functions are defined by [49]:

$$\begin{aligned}
\mathcal{R}_\Sigma(E_n) &= \frac{1}{\sqrt{\pi}} \int_0^\infty \frac{du}{\sqrt{u}} \exp(-u) \sum_{i=1}^n c_i \theta\left(\frac{u}{E_n^2} - \frac{1}{\Lambda_i^2}\right) \\
\mathcal{R}_\mathcal{Q}(E_n, E_m) &= \int_0^1 \frac{d\alpha}{2\pi} \frac{\alpha E_n - (1-\alpha) E_m}{\sqrt{\alpha(1-\alpha)}} \sum_{i=1}^n c_i \cdot \frac{\exp\left(-\frac{\alpha E_n^2 + (1-\alpha) E_m^2}{\Lambda_i^2}\right)}{\alpha E_n^2 + (1-\alpha) E_m^2} \\
\mathcal{R}_\mathcal{M}(E_n, E_m) &= \frac{1}{2} \frac{\mathrm{sign}(E_n - \mu) - \mathrm{sign}(E_m - \mu)}{E_n - E_m}.
\end{aligned} \tag{36}$$

The collective integrals reduce to [25]

$$\langle t(c,s) \mid D_{ai} \mid t(c,s) \rangle = \langle D_{ai} \rangle_c = -\frac{3}{4} S_i L_a \tag{37}$$

where S_m is the eigenvalue of the SU(2) spin generator and L_n is the eigenvalue of the left SU(3) color generators. Similarly

$$\langle t(c,s) \mid d_{ibb} D_{ab} R_b \mid t(c,s) \rangle = \langle d_{ibb} D_{ab} R_b \rangle_c = -\frac{3}{8} \;\; S_i L_a \tag{38}$$

$$\langle t(c,s) \mid D_{a8} R_i \mid t(c,s) \rangle = \langle D_{a8} R_i \rangle = \frac{\sqrt{3}}{8} \;\; S_i L_a \tag{39}$$

V. NUMERICAL RESULTS

As mentioned already, for a given top-pion decay constant f_π the constituent top mass $M_{\rm con}$ is the only free parameter of the model in the mesonic as well as in the solitonic sector. The f_π is chosen between 25 GeV and 50 GeV according to [36]. For the actual calculation we use $f_\pi = 30$ GeV. For this f_π the constituent top mass is chosen to $M_{\rm con} = 190$ GeV since then a soliton top mass around $M_{top} = 190$ GeV is achieved. Here we make the assumption that the constituent top mass $Mcon$ and the soliton top mass, M_{top}, should not differ to much. Otherwise the model and the concept of a soliton are not well defined. The corresponding scaled cutoff is then determined at $\Lambda/M_{\rm con} = 1.9$. Keeping $\Lambda/M_{\rm con}$ fixed, we varied f_π between 25 GeV and 50 GeV with, however,

negligible effect on the chromomagnetic anomalous moment κ (Tab. (III)). Hence all further and detailed calculations have been done with $f_\pi = 30$ GeV and no further adjusting of parameters has been performed. The fact that a variation of f_π has only a small effect on the anomalous chromomagnetic moment κ can be explained in the following way: Due to the dimension MeV of f_π, the influence of this constant on observables of our model can be estimated by the dimension of the observables themselves. For example the mean squared radius scales like $1/(\mathrm{MeV})^2$, so it decreases with $1/f_\pi^2$ if f_π increases. The anomalous chromomagnetic moment itself is a dimensionless quantity, because its given by the chromomagnetic moment (dimension 1/MeV) scaled by the top quark mass M_{top}. So one should expect that this anomalous moment does not depend significant on f_π, which can be seen in Tab. (III).

The final results for κ, including rotational corrections to order $\mathcal{O}(\Omega^1)$, are presented in Tab. (III) and Fig. (2). Apparently the chromomagnetic anomalous moment κ lies between 0.2 and -0.2 for a range of constituent top quark masses M_{con}, which agree roughly with the solitonic mass and are less than 350 GeV. It is interesting to investigate the effect of an explicit symmetry breaking by giving the current top quark a finite mass. The resulting effect on κ with a current mass of $m = 60$ GeV is shown for illustration in Fig. (2) as well.

Animated by Brodsky and Schlumpf [16], we plotted the anomalous magnetic moment against the product of the top soliton mass and the radius $R = \sqrt{< r^2 >}$ (see Fig. (3)). Surprisingly the result shows that $(\kappa + 1)$ is proportional to $M_{top} \cdot R$, which is supported by a linear fit of our data, shown in Fig. (3) too.

Finally we used a relation between the top production rate and κ given by Atwood, Kagan and Rizzo [20] to combine it with our relation between κ and M_{con}. So we are able to give the top production rate as a function of the top constituent mass M_{con} (Fig. (4)), the only free parameter of our model.

VI. SUMMARY

We have shown that at some high energy scale Λ the spontaneous breaking of SU(3) color not only explains the mass of the top quark in the electroweak theory but also explains an anomalous chromomagnetic moment of the top quark.

To this aim we use an effective color $SU(3) \otimes SU(3)$ invariant Lagrangian that describes the mass of the top via dynamical symmetry breaking. Soliton solutions are obtained by trivially embedding an SU(2) hedgehog Ansatz for the color pion octet into SU(3) [43] and minimizing the classical energy functional. This gives a classical mean field solution in terms of an chiral angle. The quantization of the soliton is done by semiclassical quantization of the rotational zero modes of the time independent classical solution. Imposing canonical commutation rules for the collective coordinates, generators of SU(3) color and SU(2) spin can be identified. Due to the trivial embedding the $(3, 1/2)$ triplet with spin 1/2 appears as lowest possible representation. Higher spin and color states turn out to be shifted by roughly two times the top mass.

The chromomagnetic moment turns out to be positive and of the order $-0.16 \leq \kappa \leq +0.13$ provided the constituent top quark mass is of the order of the soliton mass. Therefore it has the correct magnitude to explain the recent measured top production rate at the Tevatron [20]. To this aim we calculated the chromomagnetic moment in a solitonic background field for a pseudoscalar octet of Goldstone pions up to the linear order in the rotational velocity of the semiclassical quantization. It turned out that the recently found corrections in the linear order from explicit time-ordering of collective operators [48,50] give important contribution to the anomalous magnetic moment.

So we used a model, which is able to describe a top-quark via solitonic solutions in a quite reasonable way. If the cross section of top quark production was eventually found to agree with the SM expectations, the estimated bounds on κ – obtainable on the Tevatron – are given by $-0.14 \leq \kappa \leq 0.15$ for luminosity $\mathcal{L} = 100\,\mathrm{pb}^{-1}$ up to $-0.08 \leq \kappa \leq 0.11$ for $\mathcal{L} = 1000\,\mathrm{pb}^{-1}$, if all errors on the experimental and theoretical side are taken into account [20]. Our presented model gives an anomalous magnetic moment, which lies in the estimated bounds and also gives the new D0 and CDF results for the top cross section with a top mass in the physical magnitude.

On the other hand, if the result of D0 Collaboration [9] are correct, – and the new CDF measurements support this – there is no direct need to have an anomalous chromomagnetic moment far away from zero. However a high statistic measurement will be necessary to determine values of kappa [-0.1,0.1], which would be naturally supported in this soliton picture. But we should stress that even in the limit of vanishing kappa, our soliton model provides a picture for the top quark as a valence quark bound in a cloud of color octet Goldstone bosons. As a consequence it can have a non-vanishing radius of the order $\mathcal{O}(1am)$, contradicting the former picture of having kappa proportional to $M_{top} \cdot\ R$ [16,15].

Altogether our model may be a good instrument to investigate the new physics of the top quark in a quite simple, but effective way. The next high statistic measurements of D0 and CDF of the near future will decide, if this model can tell us more about the top.

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APPENDIX A: SU(3) MATRIX ELEMENTS

Wave functions are given by

$$\Psi^{(3)}_{(Y,T,T_3)(Y_R,J,J_3)}(A) = \sqrt{3}(-)^{Y_R/2+J_3} D^{(3)*}_{(Y,T,T_3)(-Y_R,J,-J_3)}(A) \tag{A1}$$

with $Y_R = -\frac{N_c}{3} = -\frac{1}{3}$ for convenience and $J = \frac{1}{2}$ and $J_3 = \frac{1}{2}$ for the spin 1/2 triplet. The $D^{(3)*}_{(Y,T,T_3)(-Y_R,J,-J_3)}(A)$ are the usual Wigner wave functions [51]. Then

$$\Psi^{(3)}_{(Y,T,T_3)(-1/3,1/2,1/2)}(A) = \sqrt{3}(-)^{1/3} D^{(3)*}_{(Y,T,T_3)(1/3,1/2,-1/2)}(A), \tag{A2}$$

where possible combinations for the color indices in $[\mathbf{3}]$ are

$$(Y,T,T_3) = \begin{pmatrix} 1/3 & 1/2 & 1/2 \\ 1/3 & 1/2 & -1/2 \\ -2/3 & 0 & 0 \end{pmatrix} = \begin{pmatrix} b \\ g \\ r \end{pmatrix} = color. \tag{A3}$$

Using that $[\mathbf{8}] \otimes [\bar{\mathbf{3}}] = [\bar{\mathbf{3}}] + [\mathbf{6}] + [\mathbf{15}]$ we find

$$\int d\xi_A D^{(3)}_{(Y,T,T_3)(1/3,1/2,-1/2)}(A) D^{(8)}_{ab}(A) D^{(3)*}_{(Y,T,T_3)(1/3,1/2,-1/2)}(A) =$$
$$\begin{pmatrix} 8 & \bar{3} & \bar{3} \\ a & color & color \end{pmatrix} \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ b & -\frac{1}{3},\frac{1}{2},\frac{1}{2} & -\frac{1}{3},\frac{1}{2},\frac{1}{2} \end{pmatrix},$$

where we used the relation between complex conjugation and anti-multiplets:

$$D^{(3)*}_{(L)(R)}(A) = D^{(\bar{3})}_{(-L)(-R)}(A)(-)^{Q(L)-Q(R)} \tag{A4}$$

with the charge of the states $(-R) = (-Y_R, J, -J_3)$ and $Q(R) = J_3 + Y_R/2$. Therefore in our case we have to consider

$$\langle t(c,s) \mid D^{(8)}_{33}(A) \mid t(c,s)\rangle = \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ 010 & -color & -color \end{pmatrix} \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ 010 & -\frac{1}{3},\frac{1}{2},\frac{1}{2} & -\frac{1}{3},\frac{1}{2},\frac{1}{2} \end{pmatrix}$$
$$\langle t(c,s) \mid D^{(8)}_{38}(A) \mid t(c,s)\rangle = \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ 010 & -color & -color \end{pmatrix} \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ 000 & -\frac{1}{3},\frac{1}{2},\frac{1}{2} & -\frac{1}{3},\frac{1}{2},\frac{1}{2} \end{pmatrix}$$
$$\langle t(c,s) \mid d_{3bb} D^{(8)}_{3b}(A) \mid t(c,s')\rangle = \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ 010 & -color & -color \end{pmatrix} \begin{pmatrix} 8 & \bar{3} & \bar{3} \\ b & -spin' & -spin \end{pmatrix} d_{3bb}$$

and the relation between complex conjugation and anti-multiplets:

$$D^{(3)*}_{(L)(R)}(A) = D^{(\bar{3})}_{(-L)(-R)}(A)(-)^{Q(L)-Q(R)}. \tag{A5}$$

Noting that [52,53]

$$\begin{array}{rcrcr} V_- \mid \bar{1}\rangle & = & R_{4+i5} \mid \bar{1}\rangle & = & \mid \bar{3}\rangle \\ V_+ \mid \bar{3}\rangle & = & R_{4-i5} \mid \bar{3}\rangle & = & \mid \bar{1}\rangle \\ U_- \mid \bar{1}\rangle & = & R_{6+i7} \mid \bar{1}\rangle & = & \mid \bar{2}\rangle \\ U_+ \mid \bar{2}\rangle & = & R_{6-i7} \mid \bar{2}\rangle & = & - \mid \bar{1}\rangle \\ I_- \mid \bar{2}\rangle & = & R_{1+i2} \mid \bar{2}\rangle & = & \mid \bar{3}\rangle \\ I_+ \mid \bar{3}\rangle & = & R_{1-i2} \mid \bar{3}\rangle & = & \mid \bar{2}\rangle \\ I_3 \mid \bar{3}\rangle & = & R_3 \mid \bar{3}\rangle & = & -1/2 \mid \bar{3}\rangle \\ I_3 \mid \bar{2}\rangle & = & R_3 \mid \bar{2}\rangle & = & +1/2 \mid \bar{2}\rangle \end{array} \tag{A6}$$

and rewriting $d_{3bb} D^{(8)}_{3b} R_b$ as

$$d_{3bb}D^{(8)}_{3b}R_b = \frac{\sqrt{2}}{4}\left[-D_{3p}V_+ + D_{3\Xi^-}V_- - D_{3n}U_+ - D_{3\Xi^0}U_-\right] \tag{A7}$$

gives finally, with the SU(3) Clebsch-Gordon Coefficients [55,56].

$$\begin{aligned}
\langle D^{(8)}_{33}\rangle &= -\frac{3}{4}S_3F_3 \\
\langle D^{(8)}_{38}R_3\rangle &= \frac{\sqrt{3}}{8}S_3F_3 \\
\langle d_{3bb}D^{(8)}_{3b}R_b\rangle &= -\frac{3}{8}S_3F_3
\end{aligned} \tag{A8}$$

These matrix elements are sufficient to calculate the magnetic moment up to the linear order of the rotational frequency.

TABLE I. SU(3) Clebsch Gordon Coefficients. Particle quantum numbers can be read of from Fig. 1.

$[\mathbf{8}]$	$[\bar{\mathbf{3}}]$	$[\bar{\mathbf{3}}]$	
Σ^0	$\bar{1}$	$\bar{1}$	0
Σ^0	$\bar{2}$	$\bar{2}$	$-\sqrt{\frac{3}{16}}$
Σ^0	$\bar{3}$	$\bar{3}$	$+\sqrt{\frac{3}{16}}$
Σ^+	$\bar{3}$	$\bar{2}$	$\sqrt{\frac{3}{8}}$
Σ^-	$\bar{2}$	$\bar{3}$	$-\sqrt{\frac{3}{8}}$
Λ	$\bar{1}$	$\bar{1}$	$\sqrt{\frac{1}{4}}$
Λ	$\bar{2}$	$\bar{2}$	$-\sqrt{\frac{1}{16}}$
Λ	$\bar{3}$	$\bar{3}$	$-\sqrt{\frac{1}{16}}$
Ξ^0	$\bar{1}$	$\bar{2}$	$\sqrt{\frac{3}{8}}$
Ξ^-	$\bar{1}$	$\bar{3}$	$\sqrt{\frac{3}{8}}$
p	$\bar{3}$	$\bar{1}$	$\sqrt{\frac{3}{8}}$
n	$\bar{2}$	$\bar{1}$	$-\sqrt{\frac{3}{8}}$

TABLE II. The cutoff Λ, the scaled cutoff Λ/M_{con}, the Goldstone decay constant f_π and the top condensate as a function of the top constituent quark mass M_{con}.

M_{con} [GeV]	f_π [GeV]	Λ [GeV]	Λ/M	$-\langle\bar{t}_a t_a\rangle^{1/3}$ [GeV]
185.00	30.00	368.54	1.9921	86.9087
190.00	30.00	365.49	1.9236	86.1445
200.00	30.00	360.87	1.8043	84.8325
250.00	30.00	356.09	1.4244	80.9788
350.00	30.00	382.30	1.0923	78.8805
450.00	30.00	422.20	0.9382	79.0976
550.00	30.00	466.00	0.8473	79.5127
158.30	25.0	304.59	1.9241	71.7927
190.00	30.0	365.49	1.9236	86.1455
316.70	50.0	609.12	1.9233	143.5710

TABLE III. The energy of the lowest lying top quark multiplet M_{top}, the corresponding Goldstone decay constant f_π, the moments of inertia I_1 and I_2, as well as the anomalous chromomagnetic moment κ as a function of the top constituent quark mass M_{con}.

M_{con} [GeV]	M_{top} [GeV]	f_π [GeV]	I_1 [am]	I_2 [am]	κ	$< r^2 >$ [am^2]
185	224.5(140.7)	30	1.269	.592	0.1309	2.8206
190	225.1(132.6)	30	1.114	.516	0.0923	2.4354
200	225.8(122.9)	30	.954	.440	0.0289	2.0734
250	224.3(91.3)	30	.651	.304	-0.1625	1.4183
350	215.7(35.0)	30	.458	.219	-0.3643	0.9858
450	204.3(-22.0)	30	.311	.147	-0.4625	0.7259
550	195.8(-82.8)	30	.291	.132	-0.5527	0.6561
650	186.5(-145.3)	30	.246	.111	-0.6083	0.5670
158.3	187.6(110.6)	25	1.336	.619	0.0915	3.5113
190.0	225.1(132.6)	30	1.114	.516	0.0923	2.4354
316.7	375.4(220.3)	50	.694	.319	0.1084	0.8712

TABLE IV. Same as Tab. III but with an explicit symmetry breaking current quark mass of 60 GeV.

M_{con} [GeV]	M_{top} [GeV]	f_π [GeV]	I_1 [am]	I_2 [am]	κ	$< r^2 >$ [am^2]
225	263.1(213.1)	30	2.3516	1.2034	0.1566	4.2064
230	264.7(194.3)	30	1.2967	0.5597	0.0962	2.2825
235	265.8(189.5)	30	1.1169	0.4734	0.0661	1.9879
240	266.6(185.8)	30	1.0006	0.4206	0.0400	1.8029
250	267.7(179.6)	30	0.8468	0.3543	-0.0055	1.5644

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